

Symmetries and Ambiguities in the linear sigma model with light quarks

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Abstract

We investigate the role of undetermined finite contributions generated by radiative corrections in a $SU(2) \times SU(2)$ linear sigma model with quarks. Although some of such terms can be absorbed in the renormalization procedure, one such contribution is left in the expression for the pion decay constant. This arbitrariness is eliminated by chiral symmetry.

Introduction

In the process of perturbative computations in Quantum Field Theory, the regularization scheme is a crucial ingredient for dealing with the typical divergencies of physical amplitudes. The choice of regularization is usually made at the beginning of the calculation and in this case objects like differences between integrals of the same superficial degree of divergence are automatically fixed (e.g. in dimensional regularization, Pauli-Villars regularization, etc.).

Following ideas put forward by Jackiw [1] we chose to leave undetermined until the end of calculations these *a priori* arbitrary (regularization dependent) constants which emerge in the present case from the difference of two logarithmically divergent integrals. In this context an adequate regularization scheme is Implicit Regularization (IR), proposed in [2], since all divergencies are displayed in terms of primitive divergent integrals without external momentum dependence, for which it is not necessary to explicitate a regulator. The finite integral contributions to a physical amplitude are then twofold; the ones whose integrands depend on external loop momenta and are integrated as usual and those which result from the difference of integrals with the same degree of divergence. According to renormalization theory the latter emerging arbitrary local terms correspond to the addition of finite counterterms in the Lagrangian, which may be always added as long as they comply with the underlying symmetries of the Lagrangian.

The IR method has been successfully applied in [3]-[5], corroborating and elucidating results in the literature for the case of CPT violation in extended QED, topological mass generation in 3-dimensional gauge theories, the Schwinger model in its chiral version, and also the Adler-Bardeen-Jackiw anomaly is consistently treated by the method. It has been

recently extended to massless supersymmetric theories [6]. The present work is a simple example of the use of the method in a model with strong phenomenological content. We show that although most of the arbitrary contributions which appear in the calculation can be absorbed by renormalization, chiral symmetry is required to eliminate the one which would have direct phenomenological consequences on the pion decay constant.

1 The implicit regularization technique

In the present section we illustrate the relevant technical details of IR by working out explicitly a one-loop Feynman amplitude. Consider the pseudoscalar-pseudoscalar amplitude

$$\Pi^{PP}(p^2) = i \int_{\Lambda} \frac{d^4 k}{(2\pi)^4} \text{Tr} \left\{ \gamma_5 \frac{1}{\not{k} - m} \gamma_5 \frac{1}{\not{k} - \not{p} - m} \right\} \quad (1)$$

where the symbol Λ stands for a regulator which doesn't need to be explicitated, but which is necessary to give a meaning to $\Pi^{PP}(p^2)$. Then one is allowed to manipulate the integrand algebraically. We do it in such a way that the divergencies appear as integrals and separated from the finite (external momentum dependent) contribution to eq.(1). After taking the Dirac trace, we use the following algebraic identity

$$\frac{1}{(k-p)^2 - m^2} = \frac{1}{k^2 - m^2} + \frac{2k \cdot p - p^2}{[(k-p)^2 - m^2](k^2 - m^2)} \quad (2)$$

at the level of the integrand. Note that it allows one to confine the external momentum dependence in fully convergent integrals, since this relation can be used recursively until the finite part is completely separated from divergent integrals. We get

$$\begin{aligned} \Pi^{PP}(p^2) = & -2 \left\{ i \int \frac{d^4 k}{(2\pi)^4} \frac{1}{k^2 - m^2} + i \int \frac{d^4 k}{(2\pi)^4} \frac{1}{(k-p)^2 - m^2} \right. \\ & \left. - ip^2 \int \frac{d^4 k}{(2\pi)^4} \frac{1}{(k^2 - m^2)[(k-p)^2 - m^2]} \right\}. \end{aligned} \quad (3)$$

The first integral on the RHS is what we call a basic quadratic divergence

$$I_q(m^2) = i \int \frac{d^4 k}{(2\pi)^4} \frac{1}{k^2 - m^2}. \quad (4)$$

The second integral on the RHS is also a quadratic divergence, but it still possesses an external momentum dependence. If one uses (2), one sees that an arbitrariness emerges

$$T(p^2, m^2) = i \int \frac{d^4 k}{(2\pi)^4} \frac{1}{(k-p)^2 - m^2} = I_q(m^2) + p^\mu p^\nu \Delta_{\mu\nu}, \quad (5)$$

where

$$\Delta_{\mu\nu} = i \int \frac{d^4 k}{(2\pi)^4} \frac{4k^\mu k^\nu}{(k^2 - m^2)^3} - i \int \frac{d^4 k}{(2\pi)^4} \frac{g_{\mu\nu}}{(k^2 - m^2)^2} = \alpha g_{\mu\nu}, \quad (6)$$

i.e., the difference between two logarithmically divergent integrals. In dimensional and Pauli-Villars regularizations, we would get zero for α , which is independent of the masses

in the integrals as it has been shown in [4]. The third term in eq.(3) is logarithmically divergent. Using the prescribed method we write the integral as

$$\begin{aligned} i \int \frac{d^4 k}{(2\pi)^4} \frac{1}{(k^2 - m^2)[(k - p)^2 - m^2]} &= i \int \frac{d^4 k}{(2\pi)^4} \frac{1}{(k^2 - m^2)^2} + Z_0(m^2, p^2; m^2) \\ &= i \int \frac{d^4 k}{(2\pi)^4} \frac{1}{(k^2 - \lambda^2)^2} + Z_0(m^2, p^2; \lambda^2) \end{aligned} \quad (7)$$

where again we separated another basic divergent integral

$$I_{log}(m^2) = i \int \frac{d^4 k}{(2\pi)^4} \frac{1}{(k^2 - m^2)^2}, \quad (8)$$

and a finite contribution

$$Z_0(m^2, p^2; \lambda^2) = \frac{1}{(4\pi)^2} \int_0^1 dz \ln \left(\frac{p^2 z(1 - z) - m^2}{-\lambda^2} \right). \quad (9)$$

Note that in the above expression (7) an arbitrary scale λ^2 has been introduced through the relation

$$I_{log}(m^2) = I_{log}(\lambda^2) + \frac{1}{(4\pi)^2} \ln \left(\frac{m^2}{\lambda^2} \right). \quad (10)$$

and it becomes clear that the finite term in (10) parametrizes the freedom that one has of an arbitrary constant, when separating divergent from finite contributions. The arbitrary scale λ^2 will be fixed by the choice of a renormalization point.

Finally we obtain that

$$\Pi^{PP}(p^2) = 2 \left[2I_q(m^2) + p^2 \alpha - p^2 \left(I_{log}(\lambda^2) + Z_0(m^2, p^2; \lambda^2) \right) \right]. \quad (11)$$

Analogously one finds for the scalar amplitude

$$\Pi^{SS}(p^2) = i \int_{\Lambda} \frac{d^4 k}{(2\pi)^4} \text{Tr} \left\{ \frac{1}{(\not{k} - m)(\not{k} - \not{p} - m)} \right\}, \quad (12)$$

the following expression

$$\Pi^{SS}(p^2) = 2 \left[2I_q(m^2) + p^2 \alpha + (4m^2 - p^2) \left(I_{log}(\lambda^2) + Z_0(m^2, p^2; \lambda^2) \right) \right]. \quad (13)$$

2 The model

We start with the following generating functional of the $SU(2) \times SU(2)$ linear sigma model with fermions [7] and follow closely [8]

$$Z = \int \prod_a \mathcal{D}\sigma_{0a} \mathcal{D}\pi_{0a} \mathcal{D}q \mathcal{D}\bar{q} \exp(iS(\bar{q}, q, \sigma_0, \pi_0)) \quad (14)$$

with the action

$$\begin{aligned} S(\bar{q}, q, \sigma_0, \pi_0) &= \int d^4 x \left[L_{NJL} - \frac{\beta_0}{2} \text{Tr}(\sigma_0^2 + \pi_0^2) + \frac{\mu_0^2}{2g_0} \text{Tr}(\hat{m}_0 \sigma_0) \right. \\ &\quad \left. + \frac{f_0^2}{4} \text{Tr}(\partial_\mu \sigma_0 \partial^\mu \sigma_0 + \partial_\mu \pi_0 \partial^\mu \pi_0) \right] \end{aligned} \quad (15)$$

where L_{NJL} stands for the Nambu - Jona-Lasinio Lagrangian [9] in semibosonized form [10]

$$S_{NJL} = \int d^4x \left[\bar{q} D q - \frac{\mu_0^2}{4} \text{Tr} (\sigma_0^2 + \pi_0^2) \right]. \quad (16)$$

with the Dirac operator given by

$$D = i \not{\partial} - g_0 (\sigma_0 + i \gamma_5 \pi_0). \quad (17)$$

The notation is as follows. The index 0 stands for bare quantities, the scalar σ_0 is a $SU(2)$ flavor singlet and the pseudoscalar $\pi_0 = \pi_{0i} \tau_i$ is a flavor $SU(2)$ triplet with τ_i being the usual Pauli matrices. The explicit symmetry breaking term is introduced through the term linear in the scalar field, with a factor which would correspond to the symmetry breaking pattern of the NJL Lagrangian [8]; \hat{m}_0 is a diagonal matrix with current quark masses $\hat{m}_{0u} = \hat{m}_{0d}$. To finish bosonization one integrates over the quadratic fermionic Lagrangian

$$\int \mathcal{D}q \mathcal{D}\bar{q} \exp \left(i \int d^4x \bar{q} D q \right) = \exp (\ln |\det D_E|) = \exp \left(\frac{1}{2} \text{Tr} \ln (D_E^\dagger D_E) \right), \quad (18)$$

where Tr designates functional trace and D_E stands for the euclidean Dirac operator and we chose the chiral invariant representation for the real part of the fermionic effective action [11]-[13].

As the vacuum expectation value of the scalar field acquires a finite value, one has to shift $\sigma_0 \rightarrow \sigma_0 + m/g_0$ to define the new vacuum. Here m denotes the constituent quark mass of light quarks $m = m_u = m_d$. Then we obtain

$$D_E^\dagger D_E = m^2 - \partial^2 + Y \quad (19)$$

with [14]

$$Y = i g_0 \gamma_\mu (\partial_\mu \sigma_0 + i \gamma_5 \partial_\mu \pi_0) + g_0^2 \left(\sigma_0^2 + 2 \frac{m}{g_0} \sigma_0 + \pi_0^2 \right), \quad (20)$$

leading to the expansion

$$\text{Tr} \ln D_E^\dagger D_E = \sum_{n=1}^{\infty} \frac{(-1)^{n+1}}{n} \text{Tr} \left[(-\partial_\mu^2 + m^2)^{-1} Y \right]^n. \quad (21)$$

The gap equation is obtained by considering the $n = 1$ term in eq.(21) and the linear contributions in σ_0 from the remaining terms of the bosonized action after the mass shift. We get

$$\frac{\mu_0^2}{g_0^2} (m - \hat{m}_0) - 8 N_c m I_q(m^2) + \frac{2\beta_0}{g_0^4} m^3 = 0, \quad (22)$$

where I_q has already been defined above. The $n = 2$ term in the expansion contains all the other divergent contributions which go up to four-point functions. In performing calculations there appear the two basic divergent integrals I_q and I_{log} , eqs.(4) and (8), as well as the difference between two logarithmically divergent integrals, α , eq.(5). For our purposes it suffices to focus on the radiative corrections related with the mass and kinetic terms. In the following we need the field renormalization constants. The amplitudes (11) and (13) are, up to flavor and color trace factors, the radiative corrections with two

external fields obtained from the $n = 2$ term of (21). One expands them to second order in p^2 to extract the contribution to the mesonic kinetic terms. The remaining terms of the expansion are included in the definition of the meson masses. In this way one obtains the field renormalizations for the scalar and pion fields in the form

$$Z_\sigma^{-1}(m, \lambda^2) = f_0^2 - 4N_c g_0^2 \left(I_{\log}(\lambda^2) - \alpha + Z_0(m^2, 0; \lambda^2) - 4m^2 Z'_0(m^2, 0; \lambda^2) \right), \quad (23)$$

$$Z_\pi^{-1}(m, \lambda^2) = f_0^2 - 4N_c g_0^2 \left(I_{\log}(\lambda^2) - \alpha + Z_0(m^2, 0; \lambda^2) \right), \quad (24)$$

where Z'_0 represents the derivative with respect to the external momenta squared of Z_0 , taken at $p^2 = 0$. The renormalized coupling constants $g_\sigma(m, \lambda^2)$ and $g_\pi(m, \lambda^2)$ are

$$\frac{g_0^2}{g_\sigma^2(m, \lambda^2)} = Z_\sigma^{-1}(m, \lambda^2) \quad (25)$$

$$\frac{g_0^2}{g_\pi^2(m, \lambda^2)} = Z_\pi^{-1}(m, \lambda^2). \quad (26)$$

The renormalized masses $\mu_{\sigma,\pi}(m, \lambda^2)$ become

$$\begin{aligned} \frac{\mu_{\sigma,\pi}^2(m, \lambda^2)}{Z_{\sigma,\pi}(m, \lambda^2)} &= M_{\sigma,\pi}^2(m, \lambda^2) \\ &= \mu_0^2 - 8N_c g_0^2 I_q(m^2) \\ &\quad - 4N_c g_0^2 (m \pm m)^2 \left(I_{\log}(\lambda^2) + Z_0(m^2, 0; \lambda^2) \right) \\ &\quad + 2 \frac{\beta_0}{g_0^2} (2m^2 \pm m^2) \end{aligned} \quad (27)$$

where M_σ^2 goes with the plus signs.

We obtain the renormalized quartic coupling λ_q as

$$\frac{\lambda_q}{g_\pi^2(m, \lambda^2)} = \frac{\beta_0}{g_0^4} - 4N_c I_{\log}(\lambda^2). \quad (28)$$

Note that all ambiguities α appear in the field renormalization coefficients. In principle they could assume different values for the different processes. Chiral symmetry restricts them to have the same value, as we shall see below. Finally we use the gap equation (22) in order to eliminate the quadratic divergencies in the expressions for the renormalized masses $\mu_{\sigma,\pi}^2(m, \lambda^2)$, eq.(27). It is possible to define a renormalized coupling (physical) related with the current quark mass as in [8]

$$\hat{m}_0 \frac{\mu_0^2}{g_0^2} = \frac{m \mu_\pi^2}{g_\pi^2(m, \lambda^2)}. \quad (29)$$

With these definitions we obtain the effective action

$$S'(\sigma_0 + m/g_0, \pi_0) = S_{kin} + S_{mass} + S_{int}, \quad (30)$$

with the kinetic piece

$$\begin{aligned} S_{kin} &= \int d^4x \left[Z_\sigma^{-1}(m, \lambda^2) \partial_\mu \sigma_0 \partial^\mu \sigma_0 \right. \\ &\quad \left. + Z_\pi^{-1}(m, \lambda^2) \left(\partial_\mu \pi_0^0 \partial^\mu \pi_0^0 + 2\partial_\mu \pi_0^+ \partial^\mu \pi_0^- \right) \right], \end{aligned} \quad (31)$$

and the mass terms

$$S_{mass} = - \int d^4x \left[M_\sigma^2(m, \lambda^2) \sigma_0^2 + M_\pi^2(m, \lambda^2) \left((\pi_0^0)^2 + 2\pi_0^+ \pi_0^- \right) \right]. \quad (32)$$

The interaction terms do not involve the ambiguity α , as opposed to the calculated radiative corrections with two external fields $\Pi^{PP}(p^2)$ and $\Pi^{SS}(p^2)$. This is due to the fact that after taking the Dirac trace in (21) the interaction terms are from the start at most logarithmically divergent. Recall that ambiguities emerge from the presence of quadratic divergencies with external momentum dependence at the initial stage (5). Since the interaction Lagrangian remains unchanged as compared to the one obtained in [8], we do not write it out here.

Now we display the renormalized propagators for the pions and scalar mesons

$$\Delta_\pi^{-1} = p^2 - \mu_\pi^2 + 4N_c g_\pi^2 F_{fin}(m, p^2; \lambda^2) \quad (33)$$

$$\Delta_\sigma^{-1} = p^2 - \mu_\sigma^2 + 4N_c g_\sigma^2 \Sigma_{fin}(m, p^2; \lambda^2) \quad (34)$$

with the finite momentum dependent contributions

$$F_{fin}(m, p^2; \lambda^2) = -p^2 \left[Z_0(m^2, p^2; \lambda^2) - Z_0(m^2, 0; \lambda^2) \right], \quad (35)$$

$$\begin{aligned} \Sigma_{fin}(m, p^2; \lambda^2) &= (4m^2 - p^2) \left(Z_0(m^2, p^2; \lambda^2) - Z_0(m^2, 0; \lambda^2) \right) \\ &\quad - 4m^2 p^2 Z_0'(m^2, 0; \lambda^2), \end{aligned} \quad (36)$$

where use has been made of the normalization conditions

$$\begin{aligned} \Delta_{\pi, \sigma}^{-1}(0) &= -\mu_{\pi, \sigma}^2 \\ \frac{d\Delta_{\pi, \sigma}^{-1}(p^2)}{dp^2} \Big|_{p^2=0} &= 1. \end{aligned} \quad (37)$$

At the scale $\lambda^2 = m^2$ one has $Z_0(m^2, 0; m^2) = 0$. We obtain finally the physical pseudoscalar masses as zeros of these propagators

$$\begin{aligned} m_\pi^2 &= \mu_\pi^2 - 4N_c g_\pi^2 F_{fin}(m, m_\pi^2; \lambda^2) \\ m_\sigma^2 &= \mu_\sigma^2 - 4N_c g_\sigma^2 \Sigma_{fin}(m, m_\sigma^2; \lambda^2) \end{aligned} \quad (38)$$

At the physical meson masses we obtain the following pseudoscalar quark couplings

$$\begin{aligned} g_{\pi qq}^{-2} &= g_\pi^{-2} + 4N_c \frac{dF_{fin}(m, p^2; \lambda^2)}{dp^2} \Big|_{p^2=m_\pi^2}, \\ g_{\sigma qq}^{-2} &= g_\sigma^{-2} + 4N_c \frac{d\Sigma_{fin}(m, p^2; \lambda^2)}{dp^2} \Big|_{p^2=m_\sigma^2}. \end{aligned} \quad (39)$$

3 Coupling to external fields

In order to evaluate the pion weak decay constant f_π , one introduces the axial current as an external classical field. In the present framework it appears in a generalized expression for the Dirac operator, now containing the vector V_μ and axial vector field A_μ [15, 16];

the relevant terms for the calculation of f_π are extracted after introducing the covariant derivative for the pion

$$\partial_\mu \pi \rightarrow \nabla_\mu \pi_0 = \partial_\mu \pi_0 - 2 \left(\sigma_0 + \frac{m}{g_0} \right) A_\mu - i[V_\mu, \pi_0], \quad (40)$$

in Y , eq.(20) and in the kinetic terms

$$\frac{f_0^2}{2} \text{Tr} \nabla_\mu \pi_0 \nabla^\mu \pi_0 \rightarrow -f_0^2 \frac{m}{g_0} A^\mu \partial_\mu \pi_0.$$

The expression for the pion decay constant is related to the coefficient of $p^\mu A_\mu \pi^{ph}$ where the physical pion field π^{ph} is introduced through

$$\pi_0 = \pi^{ph} \frac{g_{\pi qq}}{g_0},$$

and reads

$$f_\pi = \frac{m g_{\pi qq}}{g_0^2} \left[f_0^2 - 4N_c g_0^2 \left(I_{\log}(\lambda^2) + Z_0(m^2, m_\pi^2; \lambda^2) \right) \right]. \quad (41)$$

Note that the contribution of the radiative correction to f_π does not involve the ambiguity α , as opposed to the other radiative corrections with two external fields $\Pi^{PP}(p^2)$ and $\Pi^{SS}(p^2)$. This is again due to the fact that the Dirac trace renders the expression for the pseudoscalar-axial bubble Π_μ^{PA} logarithmically divergent

$$\Pi_\mu^{PA} = -i \int \frac{d^4 k}{(2\pi)^4} \frac{4m p_\mu}{(k^2 - m^2)[(k-p)^2 - m^2]} \quad (42)$$

Using eq.(38)

$$1 - \frac{\mu_\pi^2}{m_\pi^2} = 4N_c g_\pi^2 \left[Z_0(m^2, m_\pi^2; \lambda^2) - Z_0(m^2, 0; \lambda^2) \right] \quad (43)$$

together with eqs.(24) and (26) in eq.(41) one gets

$$f_\pi = m \frac{g_{\pi qq} \mu_\pi^2}{g_\pi^2 m_\pi^2} + \alpha' \quad (44)$$

where α' is related to α by

$$\alpha' = -4N_c m g_{\pi qq} \alpha.$$

It is now obvious that α' must be zero, given the Goldberger-Treiman relation, valid in the chiral limit

$$g_\pi f_\pi = m. \quad (45)$$

This relation is obtained from (44) with $\alpha' = 0$, since in the chiral limit $m_\pi = \mu_\pi$ as can be seen from (38) and $g_{\pi qq} = g_\pi$ from (39).

4 Conclusion

In this contribution we extended the work presented in [8] in a way as to leave open until the end of the calculations all arbitrary finite constants related with differences of divergent integrals of the same degree of divergence. This is possible using the implicit regularization technique, in which the divergent content of any amplitude is rendered strictly independent of the external (physical) momenta. It turns out that renormalization absorbs some of the arbitrary constants, while chiral symmetry fixes the arbitrary parameter appearing in the pion decay constant to be zero.

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